

Supersymmetry and gauge symmetry breaking with naturally vanishing vacuum energy

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Abstract

We review the construction of $N = 1$ supergravity models where the Higgs and super-Higgs effects are simultaneously realized, with naturally vanishing classical vacuum energy and goldstino components along gauge-non-singlet directions: this situation is likely to occur in the effective theories of realistic string models.

Invited talk presented at SUSY-95, Palaiseau, France, 15-19 May 1995

CERN-TH/95-292
November 1995

¹On leave from INFN, Sezione di Padova, Padua, Italy. Supported in part by the European Union under contract No. CHRX-CT92-0004.

1 Motivations

At the level of dimensionless couplings, the Minimal Supersymmetric Standard Model (MSSM) is more predictive than the Standard Model, since its quartic scalar couplings are related by supersymmetry to the gauge and the Yukawa couplings (for a review and references on the theoretical foundations of the MSSM see, e.g., [1]). The large amount of arbitrariness in the MSSM phenomenology is strictly related to its explicit mass parameters, the soft supersymmetry-breaking masses and the superpotential Higgs mass. Such arbitrariness cannot be removed within theories with softly broken global supersymmetry: to make progress, *spontaneous* supersymmetry breaking must be introduced.

To discuss spontaneous supersymmetry breaking in a realistic and consistent framework, gravitational interactions cannot be neglected. One is then led to $N = 1$, $d = 4$ supergravity, seen as an effective theory below the Planck scale, within which one can perform tree-level calculations and study some qualitative features of the ultraviolet-divergent one-loop quantum corrections. Of course, infrared renormalization effects can be studied, but they are plagued by the ambiguities due to the (ultraviolet) counterterms for the relevant and marginal operators. To proceed further, one must go to $N = 1$, $d = 4$ superstrings, seen as realizations of a fundamental ultraviolet-finite theory, within which quantum corrections to the low-energy effective action can be consistently taken into account, with no ambiguities due to the presence of arbitrary counterterms.

In recent years, two approaches to the problem have been followed. On the one hand, four-dimensional string models with spontaneously broken $N = 1$ local supersymmetry have been constructed [2]: none of the existing examples is fully realistic, still they represent a useful laboratory to perform explicit and unambiguous string calculations. On the other hand, many studies have been performed within string effective supergravity theories [3]: the loss in predictivity is compensated by the possibility of a more general parametrization, including possible non-perturbative effects that are still hard to handle at the string theory level. The importance of the problem and the absence of a fully satisfactory solution are reflected by the number and the diversity of the related contributions to this workshop [4, 5].

The generic problems to be solved by a satisfactory mechanism for spontaneous supersymmetry breaking can be succinctly summarized as follows.

- **Classical vacuum energy.** The potential of $N = 1$ supergravity does not have a definite sign and scales as $m_{3/2}^2 M_P^2$, where $m_{3/2}$ is the (field-dependent) gravitino mass and $M_P \equiv 1/\sqrt{8\pi G_N}$ is the Planck mass. Already at the classical level, one must arrange for the vacuum energy to be vanishingly small with respect to its natural scale.
- **$(m_{3/2}/M_P)$ hierarchy.** In a theory where the only explicit mass scale is the reference scale M_P (or the string scale M_S), one must find a convincing explanation of why the gravitino mass is at least fifteen orders of magnitude smaller than M_P (as required by a natural solution to the hierarchy problem), and not of order M_P .
- **Stability of the classical vacuum.** Even assuming that a classical vacuum with the above properties can be arranged, the leading quantum corrections to the effec-

tive potential of $N = 1$ supergravity scale again as $m_{3/2}^2 M_P^2$, too severe a destabilization of the classical vacuum to allow for a predictive low-energy effective theory.

- **Universality of squark/slepton mass terms.** Such a condition (or alternative but equally stringent ones) is phenomenologically necessary to adequately suppress flavour-changing neutral currents, but is not guaranteed in the presence of general field-dependent kinetic terms.

From the above list, it should already be clear that the generic properties of $N = 1$ supergravity are not sufficient for a satisfactory supersymmetry-breaking mechanism. Indeed, no fully satisfactory mechanism exists, but interesting possibilities arise within string effective supergravities. The best results obtained so far have been summarized in the review talk by Kounnas [5]:

- It is possible to formulate supergravity models where the classical potential is manifestly positive semi-definite, with a continuum of minima corresponding to broken supersymmetry and vanishing vacuum energy, and the gravitino mass sliding along a flat direction [6].
- This special class of supergravity models emerges naturally, as a plausible low-energy approximation, from four-dimensional string models, irrespectively of the specific dynamical mechanism that triggers supersymmetry breaking. Due to the special geometrical properties of string effective supergravities, the coefficient of the one-loop quadratic divergences in the effective theory, $\text{Str } \mathcal{M}^2$, can be written as [7, 8]

$$\text{Str } \mathcal{M}^2(z, \bar{z}) = 2Q m_{3/2}^2(z, \bar{z}), \quad (1)$$

where Q is a field-independent function, calculable from the modular weights of the different fields belonging to the effective low-energy theory. The non-trivial result is that the only field-dependence of $\text{Str } \mathcal{M}^2$ occurs via the gravitino mass. Since all supersymmetry-breaking mass splittings, including those of the massive string states not included in the effective theory, are proportional to the gravitino mass, this sets the stage for a natural cancellation of the $\mathcal{O}(m_{3/2}^2 M_P^2)$ one-loop contributions to the vacuum energy. Indeed, there are explicit string examples that exhibit this feature. If this property can persist at higher loops (an assumption so far), then the hierarchy $m_{3/2} \ll M_P$ can be induced by the logarithmic corrections due to light-particle loops.

- In this special class of supergravity models one naturally obtains, in the low-energy limit where only renormalizable interactions are kept, universal mass terms for the MSSM states ($m_0, m_{1/2}, \mu, A, B$ in the standard notation), calculable via simple algebraic formulae from the modular weights of the corresponding fields [8].

All the above results have been obtained for models where the goldstino corresponds to a gauge-singlet direction of the supergravity gauge group. In the following, we would like to summarize some recent work [9, 10] that extends the above results to models where the spontaneous breaking of supersymmetry proceeds simultaneously with the spontaneous breaking of some gauge symmetry. There are various candidates for the gauge group which

could be broken with supersymmetry (the Standard Model gauge group, some grand-unified gauge group, some hidden-sector gauge group, ...), but we do not want to be committed here to a specific realization. In our opinion, such an extension is unavoidable if one wants to incorporate the full structure of superstring models: singlet moduli of superstring effective theories are indeed charged under some gauge group broken near the string scale.

The rest of this contribution is organized as follows. In section 2 we present a toy model that illustrates some general properties of the mechanism under discussion. In section 3 we discuss two models with $SU(2) \times U(1)$ breaking. In section 4 we comment on some possible connections with SUSY GUTs, extended supergravities and four-dimensional string models.

2 A toy model

$N = 1$ supergravity models¹ are characterized by their gauge kinetic function f and Kähler function \mathcal{G} , conventionally decomposed as $\mathcal{G} = K + \log |w|^2 = -\log Y + \log |w|^2$. Consider the model, based on the Kähler manifold $[SU(1, 1)/U(1)]^3$, with

$$e^{\mathcal{G}} \equiv \frac{|w|^2}{Y} = \frac{k^2}{(S + \bar{S})(T + \bar{T})(U + \bar{U})}, \quad (k \neq 0). \quad (2)$$

It is clear that, sticking to this field parametrization, we cannot introduce any linearly realized gauge symmetry. However, by making the field redefinitions²

$$T = \frac{1 - H_1}{1 + H_1}, \quad U = \frac{1 - H_2}{1 + H_2}, \quad (3)$$

we can write

$$e^{\mathcal{G}} = \frac{k^2 |1 + H_1|^2 |1 + H_2|^2}{4 (S + \bar{S}) (1 - |H_1|^2) (1 - |H_2|^2)}. \quad (4)$$

The denominator of eq. (4) suggests two obvious $U(1)$ symmetries that can be linearly realized on the fields H_1 and H_2 , but the numerator is not invariant. However, by suitably modifying the superpotential we can move to a model described by

$$e^{\mathcal{G}} = \frac{k^2 |1 + \sqrt{H_1 H_2}|^4}{(S + \bar{S}) (1 - |H_1|^2) (1 - |H_2|^2)}, \quad (5)$$

which allows to gauge a $U(1)_X$ with charges $X(S) = 0$, $X(H_1) = -1/2$, $X(H_2) = +1/2$. Choosing for the time being a gauge kinetic function $f = S$ (this choice is not very important for the following considerations), one can observe that the Kähler metric is well-behaved in the two regions $|H_1|, |H_2| < 1$ or $|H_1|, |H_2| > 1$, that the superpotential is

¹Unless otherwise stated, we use the standard supergravity conventions where $M_P = 1$.

²Notice the danger of reasoning in terms of field VEVs and not of physical quantities: in units of M_P , the canonically normalized VEV of $(T + \bar{T})$ is always equal to 1, even when the VEV of the redefined field H_1 is equal to zero, but the two field representations correspond to the same physics.

analytic for $H_1 \neq 0$ and $H_2 \neq 0$, and that the Lagrangian is invariant under the discrete symmetries ($H_1 \rightarrow 1/H_1, H_2 \rightarrow 1/H_2$) and $H_1 \leftrightarrow H_2$. It is easy to show that the model defined above has a positive semi-definite potential ($V_F \geq 0$), and that the total classical potential $V_0 = V_F + V_D$ is minimized for arbitrary $|H_1| = |H_2|$ and S . To describe the physically inequivalent vacua, we can use the gauge-invariant VEVs $h \equiv |H_1| = |H_2|$ and $\theta \equiv \arg(H_1 H_2)$. Considering for simplicity the vacua with $\theta = 0$, and defining $s \equiv (S + \bar{S})$, the physical mass spectrum can be summarized as follows. The vector boson and gravitino masses, order parameters for gauge and supersymmetry breaking, are given by

$$m_X^2 = \frac{2h^2}{s(1-h^2)^2}, \quad m_{3/2}^2 = \frac{k^2(1+h)^2}{s(1-h)^2}. \quad (6)$$

In the spin-0 sector, there are four physical massless states, and the only massive one corresponds to $\text{Re}(H_1 - H_2)$, with mass

$$m_0^2 = m_X^2 + m_{3/2}^2 \frac{2(1+h^2)(1-h)^2}{h^2}. \quad (7)$$

In the spin-1/2 sector, there are three physical states, with masses $m_1^2 = m_{3/2}^2$ and

$$m_{2,3}^2 = m_X^2 + m_{3/2}^2 \left[1 + \frac{(1+h^2)(1-h)^2}{2h^2} \right] \pm \frac{1+h^2}{h} m_{3/2} \sqrt{m_{3/2}^2 \frac{(1-h)^4}{4h^2} + m_X^2}, \quad (8)$$

and the (canonically normalized) goldstino can be written as $\tilde{\eta} = (\hat{\tilde{S}} + \hat{\tilde{H}}_1 + \hat{\tilde{H}}_2)/\sqrt{3}$, where hats denote canonically normalized fields.

A number of observations are now in order:

- Since the goldstino components along $\hat{H}_{1,2}$ are unsuppressed, the gravitino has interactions of gauge strength via its $\pm 1/2$ helicity components [11].
- $\text{Str } \mathcal{M}^2 = -10 m_{3/2}^2$: this opens the possibility of cancelling the $\mathcal{O}(m_{3/2}^2 M_P^2)$ quantum corrections to the vacuum energy when including other sectors of the full theory. For example, n scalars with vanishing VEVs and canonical kinetic terms would give an extra positive contribution $\Delta \text{Str } \mathcal{M}^2 = 2 n m_{3/2}^2$.
- The superpotential w has a non-trivial monodromy around $h = 0$ ($H_1 \rightarrow -H_1, H_2 \rightarrow -H_2$), a situation already encountered when studying non-perturbative effects in supersymmetric theories [12].
- In the limit $m_{3/2} \ll m_X$ (which can be reached, e.g. by choosing $k \ll 1$ and h generic), the effective theory below the scale m_X would be described by $[H_S \equiv (H_1 + H_2)/\sqrt{2}]$

$$e^{\mathcal{G}} = \frac{k^2 |1 + H_S|^4}{(S + \bar{S}) (1 - |H_S|^2)^2}. \quad (9)$$

Such an effective theory would not display any singular behaviour for $h \rightarrow 0$, and would give a different value for the coefficient of the one-loop quadratic divergences, $\text{Str } \mathcal{M}^2 = -6 m_{3/2}^2$. This should remind us that a number of problems, such as

the singularity structure near the cut-off scale and the evaluation of $\mathcal{O}(m_{3/2}^2 M_P^2)$ contributions to the vacuum energy, are beyond the reach of the low-energy effective theory, and need the knowledge of the full theory to obtain meaningful answers.

- In the limit $h \rightarrow 0$, one should recover unbroken gauge symmetry ($m_X^2 \rightarrow 0$) with broken supersymmetry ($m_{3/2}^2 \rightarrow k^2/s \neq 0$), but there are some states whose masses diverge like $1/h$:

$$m_0^2 \rightarrow \frac{2m_{3/2}^2}{h^2} + \dots, \quad m_2^2 \rightarrow \frac{m_{3/2}^2}{h^2} + \dots \quad (10)$$

This is a signal that, for $h \ll \sqrt{m_{3/2} M_P}$, and denoting by Δm^2 the supersymmetry-breaking mass splittings in eq. (10), the goldstino couplings to the states in eq. (10) are of order $\Delta m^2/(m_{3/2} M_P) \sim m_{3/2} M_P/h^2 \gg 1$: this corresponds to a strongly interacting goldstino and spoils in general the reliability of perturbation theory.

- Our choice of the gauge kinetic function, $f = S$, was purely representative, and can be modified while keeping the result that $\text{Str } \mathcal{M}^2/m_{3/2}^2 = \text{constant}$. A more general form of f preserving this property is

$$f = \left(S \frac{1 - \sqrt{H_1 H_2}}{1 + \sqrt{H_1 H_2}} \right)^{-c/2} \cdot \varphi \left(S \frac{1 + \sqrt{H_1 H_2}}{1 - \sqrt{H_1 H_2}} \right), \quad (11)$$

where c is an arbitrary real constant and $\varphi(z)$ is an arbitrary holomorphic function. The original choice $f = S$ is recovered for $c = -1$ and $\varphi(z) = \sqrt{z}$. As a curiosity, observe that, choosing $\varphi(z) = z^{c/2}$, we get $f = [(1 + \sqrt{H_1 H_2})/(1 - \sqrt{H_1 H_2})]^c$. The transformation ($H_1 \rightarrow -H_1, H_2 \rightarrow -H_2$), associated with the monodromy of w around $h = 0$, would correspond in this case to a weak/strong coupling duality $f \rightarrow 1/f$.

- Another possibility is to look for different gaugings of the sigma model under consideration. For example, one could make the additional field redefinition $S = (1 - z)/(1 + z)$, and introduce the superpotential $w = k[1 + (zH_1 H_2)^{1/3}]^3$. This would allow two independent $U(1)$ factors to be gauged, producing a positive semi-definite potential, broken supersymmetry at all classical vacua, and less flat directions than in the model defined by (5). As a candidate form for the gauge kinetic function f_{ab} ($a, b = 1, 2$), it is interesting to consider in this case $f_{ab} = k_a \delta_{ab} \{ [1 + (zH_1 H_2)^{1/3}] / [1 - (zH_1 H_2)^{1/3}] \}^r$, which gives, on the vacua with $z = H_1 = H_2 \in \mathbf{R}^+$, a gaugino mass $m_{1/2} = r m_{3/2}$, and has also interesting properties with respect to weak/strong coupling duality.
- Yet another variant would consist in removing the S field (either explicitly or by introducing a superpotential that gives a VEV to its scalar component without giving a VEV to its auxiliary component), and in assigning to the fields (H_1, H_2) the Kähler potential $K = -(3/2) \log[(1 - |H_1|^2)(1 - |H_2|^2)]$ and the superpotential $w = k(1 + \sqrt{H_1 H_2})^3$. Choosing $f = L[(1 + \sqrt{H_1 H_2})/(1 - \sqrt{H_1 H_2})]^c$, with L arbitrary constant and $c \in \mathbf{R}$, would give a gaugino mass $m_{1/2} = c m_{3/2}$ at all minima with

$H_1 = H_2 \in \mathbf{R}$; the choice $c = \pm 1$ and $L \in \mathbf{R}$ would guarantee $m_{1/2}^2 = m_{3/2}^2$ at all minima, corresponding to $|H_1| = |H_2|$, but would break the discrete invariance under $(H_1 \rightarrow 1/H_1, H_2 \rightarrow 1/H_2)$.

3 $SU(2) \times U(1)$ breaking

Supergravity models of the type considered in the previous section, with gauge symmetry and $N = 1$ supersymmetry both spontaneously broken, and naturally vanishing classical vacuum energy, can be systematically constructed by generalizing the previous procedure. We would like now to discuss two examples in which the broken gauge group is $SU(2) \times U(1)$, as in the Standard Model.

A.

Consider a model based on the Kähler manifold $SU(1,1)/U(1) \times SU(2,2)/[SU(2) \times SU(2) \times U(1)]$, with the two factors parametrized by the fields S and by the 2×2 matrix

$$Z \equiv \begin{pmatrix} H_1^0 & H_2^+ \\ H_1^- & H_2^0 \end{pmatrix}, \quad (12)$$

respectively. We would like to assign, to the degrees of freedom of Z , the $SU(2) \times U(1)$ quantum numbers of the MSSM Higgs fields,

$$Z \rightarrow e^{i\alpha_A \frac{\tau^A}{2}} Z e^{i\alpha_Y \frac{\tau^3}{2}}. \quad (13)$$

In this case, choosing again $f_{ab} = \delta_{ab}S$ for simplicity, we can introduce the gauge-invariant Kähler function

$$e^{\mathcal{G}} = \frac{k^2 |1 + \sqrt{\det Z}|^4}{\det(1 - ZZ^\dagger)}. \quad (14)$$

We can easily add to the model a Kähler potential and a superpotential for the squark and slepton sectors, but we shall omit here this complication. The discussion of the model proceeds as for the toy model: inequivalent vacua are parametrized by h and θ , there are mass splittings $\Delta m^2 = \mathcal{O}(m_{3/2}^2 M_P^2/h^2)$, and h has to be chosen of the order of $G_F^{-1/2}$ to correctly reproduce the electroweak scale. This leads to a dilemma: if the gravitino mass is very light, of order h^2/M_P , then one gets mass splittings of the order of the electroweak scale, but also an unacceptable tree-level spectrum, with $\text{Str } \mathcal{M}^2 \simeq 0$ in each mass sector as in global supersymmetry; if the gravitino mass is of the order of the electroweak scale, then one gets some huge supersymmetry-breaking mass splittings, $\Delta m^2 \sim M_P^2$, and the supersymmetric solution of the hierarchy problem is endangered. Barring possible string miracles, it would seem that the gauge symmetry breaking associated with supersymmetry breaking must occur at a much heavier scale, not too far from M_P .

B.

An example along this line can be constructed with the Kähler manifold $SU(1, 1)/U(1) \times SO(2, n)/[SO(2) \times SO(n)]$, parametrized by the fields S and T, H_1, H_2, \dots , respectively, which appears for example in the effective theories of string orbifold models with tree-level supersymmetry breaking. Consider the model with $f_{ab} = \delta_{ab}S$,

$$K = -\log(S + \bar{S}) - \log[(T + \bar{T})^2 - (H_1^0 + \bar{H}_2^0)(\bar{H}_1^0 + H_2^0) - (H_1^- - \bar{H}_2^+)(\bar{H}_1^- - H_2^+) - \dots] + z^\alpha \bar{z}_\alpha, \quad (15)$$

and

$$w = k + \frac{1}{2}h_{\alpha\beta}^{(1)}z^\alpha z^\beta H_1^0 + \frac{1}{2}h_{\alpha\beta}^{(2)}z^\alpha z^\beta H_2^0 + \dots, \quad (16)$$

where the superfields z^α represent the MSSM quarks and leptons and we assume, for simplicity, the constants k , $h_{\alpha\beta}^{(1)}$ and $h_{\alpha\beta}^{(2)}$ to be real.

It is easy to show that, for $\langle z \rangle = 0$, $V_F \equiv 0$, and $V_0 = V_F + V_D$ is minimized by S, T arbitrary, $|H_1^0| = |H_2^0| \equiv h$. With the definitions $s \equiv \langle S + \bar{S} \rangle$, $t \equiv \langle T + \bar{T} \rangle$ and $x \equiv \langle H_1^0 + \bar{H}_2^0 \rangle$, the gravitino mass and the gauge boson masses read

$$m_{3/2}^2 = \frac{k^2}{s(t^2 - |x|^2)} \quad (17)$$

and

$$m_{W,Z}^2 = g_{W,Z}^2 \frac{h^2}{t^2 - |x|^2}, \quad (18)$$

respectively. Observing that, depending on the relative phase of $\langle H_1^0 \rangle$ and $\langle H_2^0 \rangle$, $0 \leq |x|^2 \leq 4h^2$, we can see that it should be $h^2/t^2 \simeq m_{W,Z}^2/M_P^2$ to reproduce correctly the electroweak scale, irrespectively of the individual values of $|x|^2$, h^2 and t^2 . In this model, all the MSSM mass terms depend on the VEVs s, t, h and x . To understand the structure of the model better, we can take the limit $h/t \rightarrow 0$, which leads to a conventional supergravity model with hidden sector and, when interactions of gravitational strength are neglected, to a special version of the MSSM. In such a limit, the MSSM mass parameters take the special values

$$\begin{aligned} m_{1/2}^2 &= m_{3/2}^2, & m_0^2(\text{matter}) &= m_{3/2}^2, & m_0^2(\text{Higgs}) &= -m_{3/2}^2, \\ \mu^2 &= m_{3/2}^2, & A^2 &= m_{3/2}^2, & B &= 0. \end{aligned} \quad (19)$$

Notice the remarkable universality properties, much more stringent than usually assumed in the general MSSM framework, with one important exception: since the kinetic terms for the Higgs and matter fields have different scaling properties with respect to the t modulus, the corresponding soft scalar masses have different values. In particular, the standard mass parameters of the classical MSSM Higgs potential are given by $m_1^2 = m_2^2 = m_3^2 = 0$, which allows for $SU(2) \times U(1)$ breaking already at the classical level, along the flat direction $|H_1^0| = |H_2^0|$.

In summary, we have seen that the structure of the toy model does not seem suitable for a direct application to $SU(2)_L \times U(1)_Y$ breaking (case A), unless one introduces some extra Standard Model singlets (case B). Indeed, we know that, in four-dimensional string models, moduli fields admit points of extended symmetry: in other words, they are charged

under some gauge group broken close to the string scale. This suggests a second intriguing possibility: to associate the breaking of supersymmetry with the breaking of a grand-unified gauge group G_U down to the MSSM gauge group. Various realizations are possible, depending on the choice of G_U and of the Kähler manifold for the Higgs sector: in this case we may perform a perturbative study of the dynamical determination of M_U and $m_{3/2}$, and there may also be applications to the doublet-triplet splitting problem of SUSY GUTs.

4 Outlook

The class of supergravity models discussed in the present talk has in our opinion rather intriguing properties (including some formal similarities with recent and less recent results on non-perturbative phenomena in globally supersymmetric theories), however it suffers from two main unsatisfactory aspects. The first is connected with the apparent arbitrariness of the construction: at the level of $N = 1$ supergravity, we are practically free to choose the gauge group, the number of chiral superfields, the Kähler manifold, the embedding of the gauge group in the isometry group of the Kähler manifold, and finally the gauge kinetic function and the superpotential that breaks supersymmetry. The second is connected with the fact that, at the level of $N = 1$ supergravity, we are essentially bound to a classical treatment, given the ambiguities of an effective, non-renormalizable theory in the control of quantum corrections, both perturbative and non-perturbative. One may hope to improve in both directions by establishing some connections with extended $N > 1$ supergravity theories and especially with four-dimensional superstring models.

To obtain a realistic $N = 1$ supergravity model, only the candidate quark and lepton superfields need to transform in chiral representations of the gauge group. It is then conceivable that the sector involved in the Higgs and super-Higgs effects can be obtained, by some suitable projection, from the gauge and gravitational sectors of an extended supergravity model. Indeed, spontaneous supersymmetry breaking with vanishing classical vacuum energy can be associated, in extended supergravities, with the gauging of a non-compact subgroup of the duality group. The examples we are aware of give gauge-singlet goldstinos in the resulting $N = 1$ theory, but one could look for models where the projected $N = 1$ goldstino transforms non-trivially under the $N = 1$ gauge group: such models would satisfy highly non-trivial constraints, due to the underlying extended supersymmetry.

Further constraints could be obtained by deriving models of the type discussed here as low-energy effective theories of four-dimensional string models with spontaneously broken $N = 1$ supersymmetry. This looks like a natural possibility: we know many examples of singlet moduli appearing in the effective string supergravities that are indeed flat directions breaking an underlying gauge group, restored only at points of extended symmetry. Unfortunately, the only existing examples are those in which supersymmetry is broken at the string tree level, via coordinate-dependent orbifold compactifications: it should be possible to study these constructions in the cases where the gauge symmetry and supersymmetry are both spontaneously broken. This could lead to some progress in the control of perturbative quantum corrections, since, working at the string level and not in the effective field theory, we can compute the full spectrum of states that contribute to the

one-loop partition function.

We hope to return to these problems in some future publication.

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